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Neutrino Magnetic Moment and the Dicyclic Group

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We show a mechanism using a discrete symmetry, the dicyclic group Q_6, which naturally gives rise to a small mass but a large transitional magnetic moment for the neutrino such that the anticorrelation between the solar neutrino intensity and the sunspot activity can be explained. The process of muonium-antimuonium oscillation provides the most interesting constraint to the model besides neutrino physics.

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Recently, the question [1] of the neutrino magnetic moment μ_ν has received a great deal of attention. A value of $\mu_\nu \sim (10^{-11}\text{-}10^{-10})\mu_B$ ($\mu_B \equiv e/2m_e$) would account for the celebrated lack of solar neutrinos [2] and furthermore would explain the apparent anticorrelation [3] between the solar neutrino flux and the sunspot activity. Namely, through μ_ν, a left-handed electron neutrino would flip [1] into its sterile right-handed component or transform [4] into the muon (or tau) antineutrino in the magnetic field of the Sun and therefore escape detection. To get a feeling for the above value of μ_ν, recall that in the standard model (SM) with ν_R included, μ_ν is enormously suppressed by the neutrino mass flip [5] $\mu_\nu(\text{SM}) \simeq 10^{-19}[m_\nu/(1\text{ eV})]\mu_B$, far below the current experimental limit [6] $\mu_\nu^e(\text{expt}) \leq 10^{-10}\mu_B$. It is not hard [7] to extend the standard model to produce a larger μ_ν, for example, by introducing a singly charged scalar boson. Then the suppression factor m_ν in the SM expression for μ_ν will be replaced by an internal fermion mass (say τ) in the extended theory, easily allowing μ_ν as large as $10^{-10}\mu_B$. The trouble, unfortunately, follows immediately since the same effect produces a large [8] neutrino mass, many orders of magnitude above the experimental limit. This can only be avoided by excessive fine-tuning which is undesirable. On the other hand, the supernova 1987A data put stringent constraints [9] on the magnetic moment of the Dirac type, but not on those of the flavor-transitional type [10].

As a way out of this, Voloshin [11] suggested using $SU(2)_\nu$ symmetry between ν_L and ν_L^c to forbid the mass term $\bar{\nu}_L^c \nu_L$, which is a triplet under the symmetry. An example adopting this idea can be found in Ref. [12]. Under such symmetry, the magnetic moment operator $\bar{\nu}_L^c \sigma_{\mu\nu} \nu_L$ is a singlet and therefore can be generated. A more realistic implementation of this $SU(2)_\nu$ symmetry would be a horizontal $SU(2)_H$ symmetry [13] between ν_{e_L} and ν_{μ_L} (in which case one can dispose of right-handed neutrinos). This $SU(2)_H$ symmetry is broken preferably in a spontaneous manner. If the $SU(2)_H$ is global, we face the problematic Goldstone boson. On the other hand, if the $SU(2)_H$ is local, the scale of symmetry breaking of $SU(2)_H$ has to be at least as large as that of the electroweak interaction, and therefore $SU(2)_H$ is lost at low energies where it is needed to avoid large m_ν. To solve this difficulty, some authors [14,15] considered the approximate $SU(2)_H$ symmetry at the Lagrangian level. We advocate the more aesthetic alternative [16–19] of using discrete symmetry, since then there is no phenomenological limit on the breaking scale.

Recently, Raffelt [20] claimed a stringent bound, $\mu_\nu < 3\times 10^{-12}\mu_B$, applicable to the transitional magnetic moment based on the core masses of red giants at the helium flash. Such a value of the magnetic moment would not be large enough to explain the solar neutrino problem. Here we take a conservative viewpoint that ignores Raffelt's bound because of the large uncertainty in our present understanding of the stellar structure. We are also aware of the Kamiokande result [21] which confirms the depletion of the solar neutrino flux but disfavors, not yet conclusively, the anticorrelation between the sunspot activity and the neutrino intensity. Even if the solar neutrino problem is due to a different explanation, our study is interesting in its own merit for suggesting the mechanism for a large neutrino magnetic moment.

In a previous paper [18], we consider a quaternion [22] symmetry Q_4, a subgroup of $SU(2)_H$ which does the same job as the whole continuous symmetry. It almost worked, in the sense that it could produce naturally small $m_\nu \lesssim 1$ eV. However, to enable $\mathbf{B}_{\text{sun}}$ to flip ν_e into ν_μ, a small fine-tuning is needed to achieve the mass difference $\Delta m_\nu^2 \equiv |m_{\nu_\mu}^2 - m_{\nu_e}^2| < 10^{-5}\text{ eV}^2$.

In this Letter we suggest a much more promising solution based on a dicyclic group [22], a twelve-element non-Abelian group which also happens to be a subgroup of $SU(2)_H$. The major advantage of having a larger group is that it possesses a Z_4 subgroup of $L_\mu - L_e$, the Zeldovich-Konopinski-Mahmoud (ZKM) symmetry [23] of the lepton number, which implies $\Delta m_\nu = 0$, as previously addressed by Leurer and Golden [24]. We will show that this Z_4 symmetry can be kept unbroken naturally. Furthermore, the model does not require any new fermion field; all of the new physics is contained in the enlarged Higgs sector.

Before we plunge into the details of our model, a discussion regarding the cosmological implication of discrete symmetries is called for. In order to allow the muon-electron mass split $m_e \neq m_\mu$, the Q_6 symmetry has to be broken. And if it is broken spontaneously, it will lead, at least according to the conventional wisdom, to the production of catastrophic domain walls when the Universe undergoes a phase transition from the unbroken phase at high temperature. A way out was emphasized by Weinberg [25] and especially by Mohapatra and Senjanović [26], through the possibility of symmetry nonrestoration in the multi-Higgs-boson systems. Even more appealing is the proposal by Linde [27], according to which a universe with a large lepton number $(n_L - n_{\bar{L}})/n_\gamma > 1$ would not undergo the usual high-temperature phase transition, even in the case of the standard model. Since the lepton number (or better a neutrino number) of the Universe is not known, this exciting scenario is still very much alive. One can also avoid this problem with a spontaneously broken discrete symmetry altogether by having the symmetry broken softly, though explicitly.

We now turn to our model, but first for the sake of completeness, discuss some essential features of Q_6 symmetry which we will use in what follows.

The dicyclic group Q_6.—Q_6 is a twelve-element non-Abelian group, a subgroup of SU(2). It is generated by elements s and r satisfying

$$s^2 = r^3 = n, \quad srs = r^2, \quad rs = s^3 r^2, \quad n^2 = e, \tag{1}$$

where e is the unit element and e,n is the center of the group. This group has four one-dimensional representations R_1, R_2, R_3, and R_4, and two two-dimensional ones, R_5 and R_6. They can be represented more explicitly by choosing a typical basis with s and r taking the values

$$\begin{aligned} &s_1 = r_1 = 1; \quad s_2 = +i, \quad r_2 = -1; \\ &s_3 = -1, \quad r_3 = 1; \quad s_4 = -i, \quad r_4 = -1; \\ &s_5 = i\sigma_3, \quad r_5 = -\exp(\tfrac{2}{3} i\pi\sigma_1); \\ &s_6 = -\sigma_3, \quad r_6 = \exp(\tfrac{2}{3} i\pi\sigma_1). \end{aligned} \tag{2}$$

Only R_5 is a faithful representation. Here we use a basis in which R_5 is represented by the SU(2) matrices. From the character table [22] of Q_6, one can write down the fusion algebra of the representations,

$$\begin{aligned} &R_5 \times R_5 = [R_1]^A + [R_3 + R_6]^S, \\ &R_6 \times R_6 = [R_3]^A + [R_1 + R_6]^S, \\ &R_5 \times R_6 = R_2 + R_4 + R_5, \\ &R_1 \times R_5 = R_3 \times R_5 = R_5, \\ &R_2 \times R_5 = R_4 \times R_5 = R_6. \end{aligned} \tag{3}$$

The superscripts A and S correspond to the antisymmetric and symmetric combinations. These fusion rules are helpful in writing down the Lagrangian with the symmetry.

The model.—Our strategy, following Voloshin, is simple: We use Q_6 to forbid the mass while allowing μ_ν. To this end, we postulate (ν_e, ν_μ) to be in the R_5 representation of Q_6 and the result follows in much the same manner as with $SU(2)_H$. Notice that we have to use R_5 because the magnetic-moment operators of Weyl neutrinos are antisymmetric in the flavor indices. Only for R_5 can the antisymmetric combination be a singlet of the group. Besides the usual doublet ϕ_s of the standard model, we shall need an additional Higgs multiplet to break the lepton-number conservation.

The particle spectrum of color singlets is displayed below, with their transformation properties under $SU(2)_L \times U(1)_Y \times Q_6$. For leptons,

$$\begin{aligned} &L_D = \begin{pmatrix} \nu_e & e \\ \nu_\mu & \mu \end{pmatrix}: \ (2, -\tfrac{1}{2}, R_5); \\ &e_R: \ (1, -1, R_4); \quad \mu_R: \ (1, -1, R_2); \\ &L_\tau = (\nu_\tau \ \ \tau): \ (2, -\tfrac{1}{2}, R_3); \quad \tau_R: \ (1, -1, R_3); \end{aligned} \tag{4}$$

for Higgs bosons,

$$\begin{aligned} &\phi_s \ (2, \tfrac{1}{2}, R_1); \quad \phi_5 \ (2, \tfrac{1}{2}, R_5); \quad \phi_6 \ (2, \tfrac{1}{2}, R_6); \\ &\eta_5^\dagger \ (1, 1, R_5); \quad \eta_6 \ (1, 0, R_6). \end{aligned} \tag{5}$$

Quarks and gauge bosons are singlets under Q_6. In (5), the subscript denotes the Q_6 representation. Furthermore, let us impose a Z_2 symmetry which transforms

$$L_D \to -L_D, \quad \phi_{5,6} \to -\phi_{5,6}, \quad \eta_{5,6} \to -\eta_{5,6}, \tag{6}$$

while the other fields are invariant. The role of these transformations will become clear from the discussion below. From (4) and (5), the Yukawa couplings are

$$\begin{aligned} \mathcal{L}_Y = {} & h_e \bar{L}_D \phi_6 e_R + h_\mu \bar{L}_D \phi_6 \mu_R + h_\tau \bar{L}_\tau \phi_s \tau_R \\ & + h_\eta L_D^T C \eta_5^\dagger L_\tau + h_5 \bar{L}_D \phi_5 \tau_R + \text{H.c.}, \end{aligned} \tag{7}$$

and the Higgs potential is

$$\mathcal{V} = m_5 (\phi_5 \phi_s) \eta_5^\dagger + m_6 (\phi_6 \phi_s^\dagger) \eta_6 + \lambda (\phi_5 \phi_6)(\eta_5^\dagger \eta_6) + \cdots, \tag{8}$$

where in Eqs. (7) and (8), Clebsch-Gordon coefficients for Q_6 are implicitly present. We display only the nontrivial terms relevant for our discussion later.

A consistent minimum of the potential is achieved with

$$\begin{aligned} &\langle \phi_s \rangle \simeq V, \quad \langle \eta_5^\dagger \rangle = \langle \phi_5 \rangle = 0, \\ &\langle \eta_6 \rangle = \langle (\eta_{6,1}, \eta_{6,2}) \rangle = (0, v), \quad \langle \phi_6 \rangle = \theta \langle \eta_6 \rangle, \end{aligned} \tag{9}$$

where we write vacuum expectation values (VEV's) in the Q_6 space and θ ($\lesssim 1$) is a mixing angle to be discussed. The electroweak scale is set by V. Equation (9) implies that even after Q_6 is broken by $\langle\phi_6\rangle$ and $\langle\eta_6\rangle$, its Z_4 subgroup generated by s remains unbroken since $s\langle\phi_6\rangle=\langle\phi_6\rangle$ and $s\langle\eta_6\rangle=\langle\eta_6\rangle$. This Z_4 symmetry also protects the naturalness of the vanishing VEV's of $\phi_{6,1}$, $\eta_{6,1}$, ϕ_5, and $\eta_5^\dagger$ in (9). From (7), the τ lepton gets the mass from the large VEV $\langle\phi_s\rangle$, whereas the electron and the muon are naturally lighter, since

$$m_e \sim \theta h_e v, \quad m_\mu \sim \theta h_\mu v, \quad m_\tau \sim h_5 V. \tag{10}$$

In order to keep the Q_6 symmetry to as low energies as possible, we assume $v\sim 1$ GeV, $\theta \sim 0.1$. Notice that the model offers no explanation of why $m_e \ll m_\mu$ (just $h_\mu \simeq 200 h_e$).

A technical remark is in order. One may worry about the nonzero $\langle\phi_6\rangle$, since if it is driven by a conventional negative mass-squared term, the masses of the two neutral Higgs bosons will be around a few GeV naturally. Phenomenologically, the Z gauge boson will decay into these Higgs bosons, with the partial width equivalent to that of one light ν generation. This is ruled out [28] experimentally. For this reason, we introduced the $SU(2)_L \times U(1)_Y$ singlet η_6 which is allowed to have a negative mass-squared term. Through the m_6 term in (9), it will drive a nonzero VEV for ϕ_6 even with a positive mass-squared term $m_{\phi_6}^2$ for ϕ_6. Therefore $\langle\phi_6\rangle \simeq m_6\langle\phi_s\rangle\langle\eta_6\rangle/m_{\phi_6}^2 < \langle\eta_6\rangle$ for $m_6 \lesssim m_{\phi_6} \sim M_W$. All we need is a small, but not tiny, mixing angle, $\theta = \langle\phi_6\rangle/\langle\eta_6\rangle \sim 0.1$, which controls the partial width of the Z^0 decay.

Before we proceed with the determination of the magnetic moments and masses, one last technical comment is called for. The Z_4 unbroken symmetry of Q_6 generated by s is

$$\begin{aligned} &e_{L,R} \to -ie_{L,R}, \quad \mu_{L,R} \to i\mu_{L,R}, \quad \tau_{L,R} \to -\tau_{L,R}, \\ &\nu_e \to -i\nu_e, \quad \nu_\mu \to i\nu_\mu, \quad \nu_\tau \to -\nu_\tau, \end{aligned} \tag{11}$$

which means that ν_τ is decoupled from ν_e and ν_μ in both the mass terms and the magnetic-moment interaction. In particular, only the transitional moment $\bar{\nu}_e^c \sigma_{\alpha\beta}\nu_\mu$ exists and only the $\nu_e\nu_\mu$ mass term is allowed. Simply, $s = i(L_\mu - L_e)$.

Magnetic moment.—The leading contribution to Q_6-invariant μ_ν is given in Fig. 1:

$$\mu_\nu \simeq e m_\tau \frac{h_5 h_\eta}{16\pi^2} \frac{m_5\langle\phi_s\rangle}{m_{\eta_5}^2 m_{\phi_5}^2}. \tag{12}$$

The diagram illustrates clearly the need for $\eta_5^\dagger$ and ϕ_5. The insertion of $m_5\langle\phi_5\rangle$ into the diagram gives rise to the Q_6-conserving mixing between the two doublets (R_5's), ϕ_5 and η_5. For generic values of $m_5 \sim m_{\eta_5} \sim m_\phi \sim M_W$, $h_5 \sim h_\eta \sim g$, it gives

$$\mu_\nu \simeq G_F m_e m_\tau \mu_B/\pi^2 \simeq (10^{-10}\text{–}10^{-11})\mu_B. \tag{13}$$

As we stressed in the introduction, a large μ_ν can be ob-

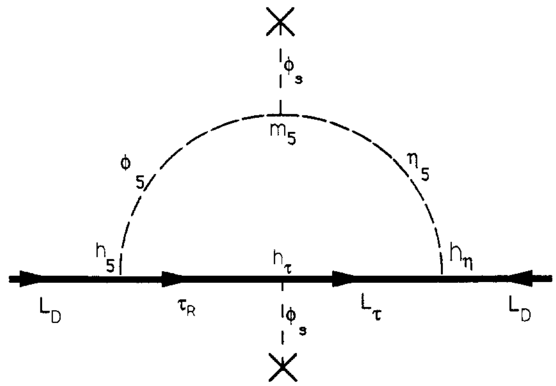

FIG. 1. A typical diagram for the transitional magnetic moment of the neutrino. The photon line is attached to any internal charged lines. There is a similar diagram for the neutrino mass with the removal of the photon line and the replacement of the $m_5\langle\phi_5\rangle$ insertion by the $\lambda\langle\eta_6\rangle\langle\phi_6\rangle$ insertion. Alternatively, one can keep $m_5\langle\phi_5\rangle$ insertion and attach additional insertions with Q_6-violating VEV on the Higgs-boson propagators to give the neutrino mass.

tained naturally. The anticorrelation of the solar neutrino and the sunspot activity demands the value of μ_ν in (13), which can be achieved using natural values of various parameters of the theory.

Neutrino mass.—As a result of (13), we only have to worry about the neutrino mass term $\nu_e^T C \nu_\mu$. The typical contribution is generated at the one-loop level from a modification of the diagram of Fig. 1, by removing the photon line and replacing the $m_5\langle\phi_5\rangle$ insertion by the $\lambda\langle\eta_6\rangle\langle\phi_6\rangle$ insertion. The new insertion gives rise to the Q_6-violating mass differences between the two components of the Q_6 doublets, ϕ_5 and $\eta_5^\dagger$. The estimate gives

$$m_{\nu_e\nu_\mu} \simeq \lambda m_\tau \frac{h_5 h_\eta}{16\pi^2} \frac{v^2}{m_H^2}\theta, \tag{14}$$

where m_H is the largest Higgs-boson mass in the loop. For generic values, $\lambda \sim e$ and $m_H \sim M_W$, it gives $m_{\nu_e\nu_\mu} \simeq \mu_\nu v^2\theta$. Here lies the reason for the Z_2 symmetry of (6); it provides additional quadratic suppression in v. For $\mu_\nu \sim 10^{-11}\mu_B$, $m_{\nu_e\nu_\mu} \simeq \theta \times (3 \text{ eV})$, well below the current experimental bounds. There are other one-loop diagrams with Q_6-violating insertions that can contribute to $m_{\nu_e\nu_\mu}$. However, they are all of the same order of magnitude as that in Eq. (14) and therefore our analysis above is not affected.

Discussion.—(a) The essential condition of ν_e-ν_μ flip in the magnetic field of the Sun is that the mass difference between ν_e and ν_μ be small: $\Delta m_\nu \lesssim 10^{-2}\text{–}10^{-3}$ eV. Because of the Z_4 subgroup of $L_\mu - L_e$, in our case, the neutrino mass matrix is of the Dirac type which implies $\Delta m_\nu = 0$ exactly. However, some numerical analyses [4] give an indication that a resonant effect which mimics the Mikheyev-Smirnov-Wolfenstein phenomenon [29] may be necessary to get large enough oscillations. If that is the case, one needs a small, but nonzero, Δm_ν of order $10^{-2}\text{–}10^{-3}$ eV. This can presumably be achieved by a

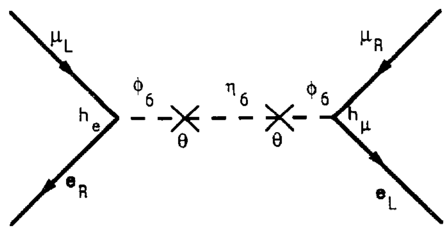

FIG. 2. A diagram for the muonium-antimuonium conversion.

small breaking of the Z_4 symmetry.

(b) Again, due to Z_4, processes like $\mu \to e\gamma$ and $\mu \to ee\bar{e}$ are strictly forbidden. The most interesting constraint on the model comes from the muonium-antimuonium oscillation [30], $(\mu\bar{e})$-$(\bar{\mu}e)$, consistent with the full $L_\mu - L_e$ symmetry. The estimate for the strength of the effective four-fermion vertex $\bar{e}\gamma_\alpha(1-\gamma_5)\mu\bar{e}\gamma^\alpha(1+\gamma_5)\mu$ (see Fig. 2) is $(\theta^2/8)(m_e m_\mu/\langle\phi_6\rangle^2)M_{\text{light}}^{-2}$. The experimental bound [31] is about $0.4G_F$, which is satisfied with our assumption that the mass of the light boson associated with η_6 and ϕ_6 be about $v \sim 1$ GeV.

(c) The new charged scalar bosons from ϕ_6 modify the $V-A$ nature of $\mu \to e\bar{\nu}_e\nu_\mu$ decay and imply the limit $h_e h_\mu < 10^{-2}$. This is in accord with $m_e m_\mu \simeq h_e h_\mu \langle\phi_6\rangle^2$ or $h_e h_\mu \simeq 5\times 10^{-3}$.

After finishing this manuscript, we became aware of a paper by Ecker and Grimus [32] which discusses the role of ZKM symmetry associated with some horizontal symmetries in the issue of the neutrino magnetic moment.

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